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Light-Induced Drift of Quantum-Confined Electrons in Semiconductor Heterostructures

by

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# LIGHT-INDUCED DRIFT OF QUANTUM-CONFINED ELECTRONS IN SEMICONDUCTOR HETEROSTRUCTURES

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The effect of light-induced drift of quantum-confined electrons in semiconductor heterostructures is predicted. The effect manifests itself as the electric current in the heterostructure plane in response to optical excitation with the frequency close, but not exactly equal, to a transition frequency between levels in the heterostructure. The current reverses its direction with a change in the detuning sign, and vanishes if the radiation polarization is normal to the heterostructure plane.

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The effect of light-induced drift (LID) in gases, predicted by Gel'mukhanov and Shalagin<sup>1</sup> and observed later<sup>2</sup>, is manifested in mechanical drift of the absorbing atoms (molecules) in response to optical excitation. The effect of LID takes place if the atom interaction with light is velocity-dependent (due to the Doppler effect), the spectrum of the exciting radiation is asymmetric with respect to the absorption line, and the absorbing molecules experience collisions with a buffer gas, with the collision rate dependent upon their internal state.

The idea of LID is the following. Suppose for the sake of definiteness that the collision frequency of the absorbing molecules with the buffer is greater in the excited state and that the exciting light is red-shifted with respect to the absorption contour. Then due to the Doppler shift, the exciting radiation interacts mainly with the counter-moving molecules, which undergo transitions into an excited state, in which their friction upon the buffer component is increased. Due to this increase, the counter-flying molecules are slowed down, while the molecules moving in the direction of light are not affected. The net effect is a drift of the absorbing component in the direction of light. The drift alternates its direction with the change of signs of either the detuning or the difference in the collision frequencies.

Later, a similar effect, surface LID, was predicted<sup>3</sup> and observed<sup>4</sup> to take place in a one-component gas whose scattering from the walls of the container is dependent on the molecule internal state. Also, it has been predicted<sup>5,6</sup> and observed<sup>7</sup> that there exists LID of electrons in semiconductors in the presence of a magnetic field. In this case the translational (Doppler-shifting) motion occurs along the field direction, and light excites transitions between Landau levels ("internal" states of the electron). If the electron translational relaxation depends on which of the Landau levels the electron occupies, LID takes place.

In this Letter we predict LID for quantum-confined carriers in semiconductor heterostructures: quantum wells, wires, and one- and two-dimensional superlattices. The heterostructures discussed are of great interest from the point of view of various applications in microelectronics and quantum optics (see, e.g., Ref. 8). The effect of LID manifests itself as the parallel (with respect to the heterostructure) current of the carriers with an

antisymmetric dependence on the detuning of the exciting radiation from the transition between states in the well, which is characteristic of LID.

For definiteness, let us consider electrons in quantum wells. Straightforward generalization of the theory is possible to describe quantum wires as confined structures and also periodic quantum structures (superlattices). The quantum well is a flat layer in the semiconductor with different chemical composition. Electrons are confined within this layer, which can be considered as a one-dimensional (say, along the  $z$ -coordinate) potential well, while their (quasi)momentum  $p$  in the  $xy$ -plane is conserved. Light causes transitions between states (subbands) in the well (so-called QWEST). The electrons in the quantum well can be considered as a two-dimensional gas of particles, whose internal state is the state in the well. The electron translational relaxation can depend on this internal state (see below). If so, the conditions of the existence of LID are met. Note that a similar idea concerning thin films and inversion layers was mentioned in Ref. 6 but was not examined there.

The LID velocity  $v_d$  can be estimated, and the necessity of the quantum confinement can be understood from the following estimates. Let us consider the transitions between two electron bands  $m$  and  $n$ , with energies at origin  $\varepsilon_a$  and dispersion laws  $\varepsilon_p^{(a)}$ , where  $a = m, n$ . The resonance condition is  $\varepsilon_n - \varepsilon_{p+q}^{(n)} - \varepsilon_m - \varepsilon_p^{(m)} - \hbar\omega \lesssim \hbar\Gamma$ , where  $q$  and  $\omega$  are the photon momentum and energy, and  $\Gamma$  is the transition homogeneous width. This condition should be met for a resonant momentum  $p_0$ , but not for momenta  $p$  perpendicular to  $p_0$ , which ensures the selectivity of excitation in the electron velocities and, consequently, existence of the drift. Let us suppose non-degenerate electron gas and use the approximation  $\varepsilon_p^{(a)} = p^2/2m_a^*$ , where  $m_a^*$  is the corresponding effective mass. The magnitude of the drift velocity  $v_d$  is proportional to and limited by the thermal velocities of the electrons participating in the resonant transitions  $v_0 \sim p_0/m^*$ . The maximum drift velocity,  $v_d \sim v_0$ , is achieved when the relaxation rates in the electronic states  $m$  and  $n$  differ strongly. Since LID is based on the thermal movement, the magnitude of  $v_d$  is increasing with the temperature  $T$ , unless  $T$  (in energy units) becomes greater than the energy gap  $\varepsilon_{nm} = \varepsilon_n - \varepsilon_m$ , where the population of the two bands becomes equal and the optical absorption vanishes. In accord with this, we shall put  $T \sim \varepsilon_{nm}$ .

For properties of LID it is of principal importance that  $q$  is much smaller than the characteristic thermal momentum  $p_t \sim m^* v_t$ . The ratio  $q/p_t \sim (\varepsilon_{nm}/m^* c^2)^{1/2}$ , where  $c$  is the velocity of light, can be estimated for the parameters typical for GaAs/AlGaAs wells.  $m^* = 0.068m_e$ , where  $m_e$  is the electron mass, and  $\varepsilon_{nm} = 30$  meV, as  $q/p_t \sim 10^{-3}$ . Taking into account this small parameter, from the analysis of the above conditions for parallel bands,  $m_n^* = m_m^*$ , which is the case for the intersubband transitions, one gets  $v_0 \sim v_t$ , where  $v_t \sim (\varepsilon_{nm}/m^*)^{1/2}$  is the mean thermal velocity of electrons. In the case of nonparallel bands (interband transitions),  $v_0 \sim v_t(m^*/\Delta m^*)(q/p)$ , where  $\Delta m^* = m_n^* - m_m^*$ , i.e. the maximum drift velocity is much smaller. From this result it is also clear that the condition of the parallel bands is  $\Delta m^* \lesssim q/p_t$ . Note that the maximum drift velocity induced by photon-drag (light-pressure) can be estimated as  $q/m^*$  and is by the factor  $q/p_t$  less than the LID velocity in the case of parallel bands. Based on the above estimates, we shall now consider the intersubband (QWEST) transitions of the quantum-confined electrons. This ensures parallel bands and maximum drift velocity.

Assuming low electron density and neglecting interaction between electrons, the Hamiltonian of the system has the form

$$H = \sum_{pa} (\varepsilon_a - \varepsilon_p) \alpha_{ap}^\dagger \alpha_{ap} + \sum_{pqab} V_{ab}(\mathbf{q}) \alpha_{ap-\mathbf{q}}^\dagger \alpha_{bp} \quad (1)$$

Here  $\alpha^\dagger$  and  $\alpha$  are electron creation and annihilation operators;  $a, b$  denote states in the well;  $p, q$  are the two-dimensional momenta of electron and photon in the  $xy$ -plane; and  $V_{ab}$  is the matrix element of the interaction with light.  $V_{ab}(\mathbf{q}) = \hbar G_{ab} \exp(i\omega t)$ , where  $G_{ab} = -e\mathbf{E}_q \mathbf{z}_{ab}/\hbar$  is the Rabi frequency,  $e$  is the carrier charge,  $\mathbf{E}_q$  is the electric field of the light with the wave vector  $\mathbf{k} = \mathbf{q}/\hbar$ , and  $\mathbf{z}$  is the normal component of the particle coordinate vector.

We make use of the one-electron density matrix in the Wigner representation.  $\rho_{ab}(\mathbf{p}, \mathbf{r}) = \sum_{\mathbf{q}} \exp(-i\mathbf{q}\mathbf{r}/\hbar) \langle \alpha_{a\mathbf{p}+\frac{\mathbf{q}}{2}}^\dagger \alpha_{b\mathbf{p}-\frac{\mathbf{q}}{2}} \rangle$ , where  $\mathbf{r}$  is the electron  $xy$ -coordinate vector. The equations of motion for  $\rho$  are obtained in the usual way by commuting  $\alpha_{a\mathbf{p}+\frac{\mathbf{q}}{2}}^\dagger \alpha_{b\mathbf{p}-\frac{\mathbf{q}}{2}}$  with the Hamiltonian (1) using the Fermi commutation rules. Adding the relaxation  $R$  and scattering  $S$  terms, and taking into account that  $q \ll p$ , one obtains the quantum

Liouville equation

$$\frac{\partial \rho}{\partial t} + \mathbf{v} \frac{\partial \rho}{\partial \mathbf{r}} = \frac{i}{\hbar} [\rho, \varepsilon + V] - R + S \quad (2)$$

where  $\varepsilon$  and  $V$  are one-electron operators with the matrix elements  $(\varepsilon)_{ab} = \varepsilon_n \delta_{ab}$  and  $V_{ab}(\mathbf{r})$ , and  $\mathbf{v} = \partial \varepsilon_p / \partial \mathbf{p}$  is the electron group velocity. We will consider only two combining states  $m$  and  $n$ . For the relaxation rate of  $\rho_{mn}$  we take the form  $(R-S)_{mn} = -\Gamma \rho_{mn}$ , where  $\Gamma$  is the polarization relaxation constant. The population decay rates shall be described in the strong collision model, in which  $S_{aa} = -\nu_a(\rho_a - \rho_a^{(0)})$ , where  $\rho_a^{(0)}$  is the equilibrium population of the state  $a = m, n$ , and  $\nu_a$  is the collision frequency for the electron in this state. The noncollisional (e.g. radiational) decay rate  $R_{aa}$  will be neglected for simplicity.

For the sake of the present Letter, it is sufficient to find the solution  $\rho_{ab}$  within perturbation theory in the field  $E$  starting with the zeroth-order  $\rho_a^{(0)} \delta_{ab}$  in the conventional Fermi-distribution form. The (two-dimensional) current density  $\mathbf{j}$  is expressed in terms of the second-order density matrix  $\rho^{(2)}$ ,

$$\mathbf{j} = e \int \mathbf{v} [\rho_{mm}^{(2)}(\mathbf{p}) + \rho_{nn}^{(2)}(\mathbf{p})] \frac{d\mathbf{p}}{(2\pi\hbar)^2} \quad (3)$$

For an isotropic dispersion law ( $\varepsilon_p = \varepsilon_p$ ), the current is directed along  $\mathbf{k}$ , and its magnitude with the use of solution of Eq. (2) can be presented in the form

$$j = I \frac{8\pi e^3}{\hbar^2 c} e_z \varepsilon_{mn}^2 \Gamma (\nu_n^{-1} - \nu_m^{-1}) \int_{-\infty}^{\infty} v [\rho_n^{(0)}(p) - \rho_m^{(0)}(p)] [(\omega - \omega_{nm} - kv)^2 - \Gamma^2]^{-1} dp \quad (4)$$

where  $I$  is the light intensity,  $\omega_{nm} = \varepsilon_{nm}/\hbar$ , and  $e_z$  is the  $z$ -component of the light polarization vector.

For simplicity, returning to the nondegenerate electron gas in the well and quadratic dispersion law, from (4) we obtain the current density in the form

$$j = I e_z \varepsilon_{mn}^2 \frac{8\pi^{3/2} e^3}{\hbar^2 c} \frac{z_{mn}^2 n}{kZ} (\nu_n^{-1} - \nu_m^{-1}) [\exp(-\varepsilon_n/T) - \exp(-\varepsilon_m/T)] C(\xi) \quad (5)$$

where  $n$  is two-dimensional electron density in the well,  $Z = \sum_a \exp(-\varepsilon_a/T)$  is the statistical sum, and the real function  $C$  and its complex argument  $\xi$  are defined as  $C(\xi) \equiv \text{Re} \left\{ \xi \exp(-\xi^2) [1 - \text{erf}(i\xi)] \right\}$ ,  $\xi \equiv \xi' - i\xi''$ ,  $\xi' = (\omega - \omega_{nm})/kv_t$ ,  $\xi'' = \Gamma/kv_t$  and  $v_t = (2T/m^*)^{1/2}$ .

The current density  $j$  (4) has characteristic polarization dependence: it vanishes for the polarization parallel to the well plane. The spectral contour of  $j$  is determined by the function  $C(\xi)$ , which is plotted in Fig. 1 using a convenient normalization of both abscissa and ordinate values. This contour has a typical for LID antisymmetric spectral dependence (the variable  $\xi'/\xi''$  is simply the detuning in the units of the homogeneous width  $\Gamma$ ). The current  $j$  increases with the intensity  $I$  leveling off in the saturation regime. As a function of  $T$  the current (5) is reaching its maximum at  $T \sim \varepsilon_{nm}$ , provided  $\Gamma = \text{const}$  in  $T$ . Though Eq. (5) is valid only for small  $I$ , the maximum (saturated) current  $j_{\text{max}}$  can easily be estimated.

Consider first the most favorable case of large Doppler broadening ( $\xi'' \ll 1$ ), where the width of the absorption contour is  $kv_t$ , with the maximum value of  $C(\xi)$  on order of unity. Assuming the collisions to determine the electron lifetime, the saturation condition is  $G_{mn}^2/kv_t \sim \min(\nu_m, \nu_n)$ . Taking this all into account, we obtain from Eq. (5) the simple estimate

$$j_{\text{max}} \sim env_t(\nu_m - \nu_n) \left[ \max(\nu_m, \nu_n) \right]^{-1} Z^{-1} \left[ \exp(-\varepsilon_n/T) - \exp(-\varepsilon_m/T) \right]. \quad (6)$$

As one can see from Eq. (6), the maximum drift velocity  $v_d$  can be comparable to the thermal velocity of electrons.

In the opposite limiting case where the homogeneous broadening is large compared to the Doppler broadening ( $\xi'' \gg 1$ ), the saturation condition is  $G_{mn}^2, \Gamma \sim \min(\nu_m, \nu_n)$ . In this case, we get from Eq. (5) the estimate of the current density

$$j_{\text{max}} \sim enkv_t^2 \Gamma^{-1} (\nu_m - \nu_n) \left[ \max(\nu_m, \nu_n) \right]^{-1} Z^{-1} \left[ \exp(-\varepsilon_n/T) - \exp(-\varepsilon_m/T) \right]. \quad (7)$$

As we can see, in this case the effect is lower as compared to Eq. (6) by a factor of  $kv_t/\Gamma$ . This is understandable since LID is based on Doppler broadening.

Let us discuss the physical situation in regard to an experimental examination of the effect suggested. We assume, as discussed above, intersubband (QWEST) transitions in GaAs/AlGaAs quantum wells,  $T \sim \varepsilon_{nm} \approx 30$  meV, the latter corresponding to the well depth  $l \approx 140$  Å. To have appreciable concentration  $n$  of electrons in the conduction band inside the well, barrier or well regions should be doped with donor impurities. These



impurities may simultaneously serve as the scatterers needed for LID (see also below). The usually achieved concentrations are  $n = 10^9$  to  $10^{12} \text{ cm}^{-2}$ . At these concentrations, electron-electron scattering does not play a significant role in relaxation. Let us also note that the electron concentration may also be controlled by normal electric field or incoherent optical pumping. The energies and temperatures assumed are below the spectral gap of optical phonons, which are therefore unlikely to contribute considerably to the electron relaxation. The role of acoustic phonons for the processes under consideration is not very important. Thus, under the present conditions, the main origin of the polarization relaxation is the broadening of the transition due to the variation  $\Delta l$  of plus or minus one atomic monolayer in the depth  $l$  of the well. This brings about the relative width  $\hbar\Gamma/\varepsilon_{nm} \approx 2\Delta l/l \approx 0.04$ , which is adopted below. This value is in fairly good agreement with the line widths experimentally observed<sup>9,10</sup> at both 60 K and 2 K (the absence of a temperature effect confirms the non-phonon broadening mechanism).

In the case under consideration, the dependence  $j_{maz}(T)$  is that explicitly shown in Eqs. (7), and it levels off for  $T \gtrsim \varepsilon_{nm}$ , i.e. for the temperatures in the liquid nitrogen range. Adopting for definiteness  $T = 77 \text{ K}$ , we get  $kv_t/\Gamma \approx 0.015$ . This corresponds to the case of weak Doppler broadening, so that Eq. (5) is to be used, and the drift velocity  $v_d$  constitutes a few percent of the thermal velocity. Assuming a large difference in the collision frequencies and, as an example, adopting the width  $L$  of the well (in the  $xy$ -plane) to be 1 mm, from (5) for  $n \sim 10^9$  to  $10^{12}$  we obtain from (7) the estimate for the current  $J_{maz} = Lj_{maz} \sim 1 \text{ } \mu\text{A}$  to 1 mA, which is more than sufficient to observe the effect and also promises the possibility of applications.

If the ends of the specimen are not electrically connected, as was assumed above, then LID brings about the appearance of the potential difference between these ends,  $\Delta U = jL_k/(\mu n e l)$ , where  $j$  is given by Eq. (5),  $L_k$  is the length of the specimen in the  $k$ -direction, and  $\mu$  is the electron mobility. In practice, the voltage may be sensitive to shunting of the well by the bulk conductivity of the specimen and is, therefore, a less suitable characteristic to measure.

The difference  $\nu_m - \nu_n$  originates from the fact that the distributions of both scatterers and confined electrons depend upon  $z$ , with the latter being state dependent.

To illustrate this, consider a realistic model in which the translational relaxation is due to scattering by a potential of impurities with concentrations  $n_i$  and  $n_o$  inside and outside the well. Applying Fermi's golden rule, and assuming that scattering from impurities is inelastic (i.e. causing the transitions  $m \rightarrow n$ ) and the well is rectangular, we obtain for the relaxation frequencies the expression with clear physical meaning

$$\nu_a = f [(n_i - n_o)P_a - n_o] \quad . \quad P_a = \int_0^l \phi_a(z)^2 dz \quad . \quad (8)$$

where  $P_a$  is the probability for an electron in the state  $a = m, n$  to be inside the well ( $\phi_a$  is the corresponding wave function) and  $f$  is a state-independent parameter determined by the interaction. The dependence of  $P_a$  upon the state  $a$  is illustrated in Fig. 2, where  $P_a$  is plotted as a function of the dimensionless parameter  $u = U_0 m^* l^2 / \hbar^2$ , with  $U_0$  being the potential well depth, for the three lowest levels. For the well considered,  $u \approx 20$ , so that  $P_a \approx 1$ , and the most favorable condition for LID is  $n_o \gg n_i$ .

To summarize, we predict a new effect, LID of electrons in quantum-confined semiconductor heterostructures. This effect differs from photon-drag and some other optoelectric phenomena in semiconductors by its magnitude and the characteristic spectral and polarization dependences.

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## FIGURE CAPTIONS

- Fig. 1. Spectral factor  $\xi''^2 C(\xi)$  as a function of  $\xi'/\xi''$  with  $\xi'$  as a variable and the values of  $\xi'' = 0.5$  (dot-dashed curve), 2.0 (dashed curve) and 50 (solid curve).
- Fig. 2. Probability  $P_a$  to be inside the well as a function of the parameter  $u$ . The solid, dashed and dot-dashed curves correspond to the three lowest levels in the well ( $a = 0, 1, 2$ ).

Fig. 1

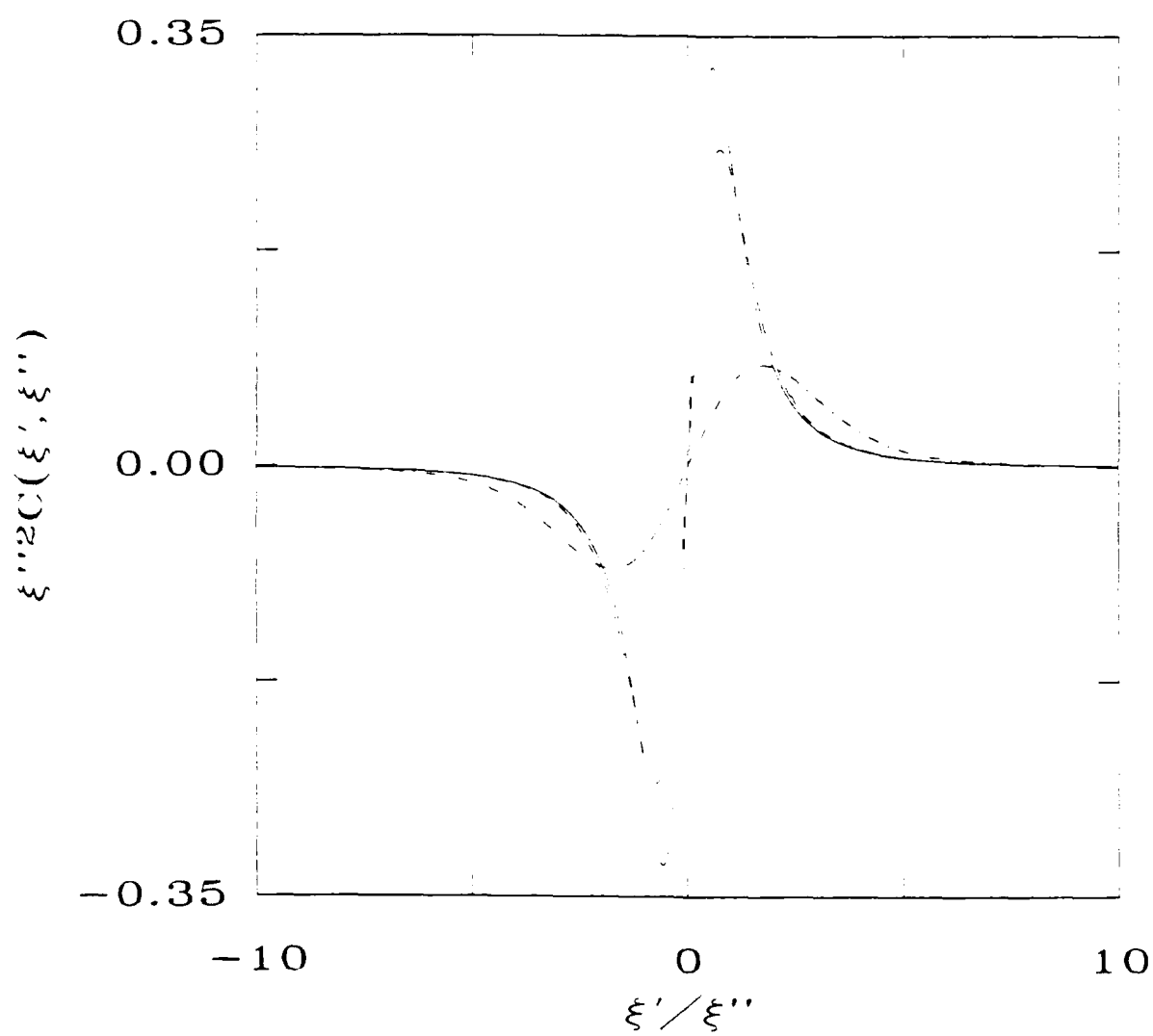
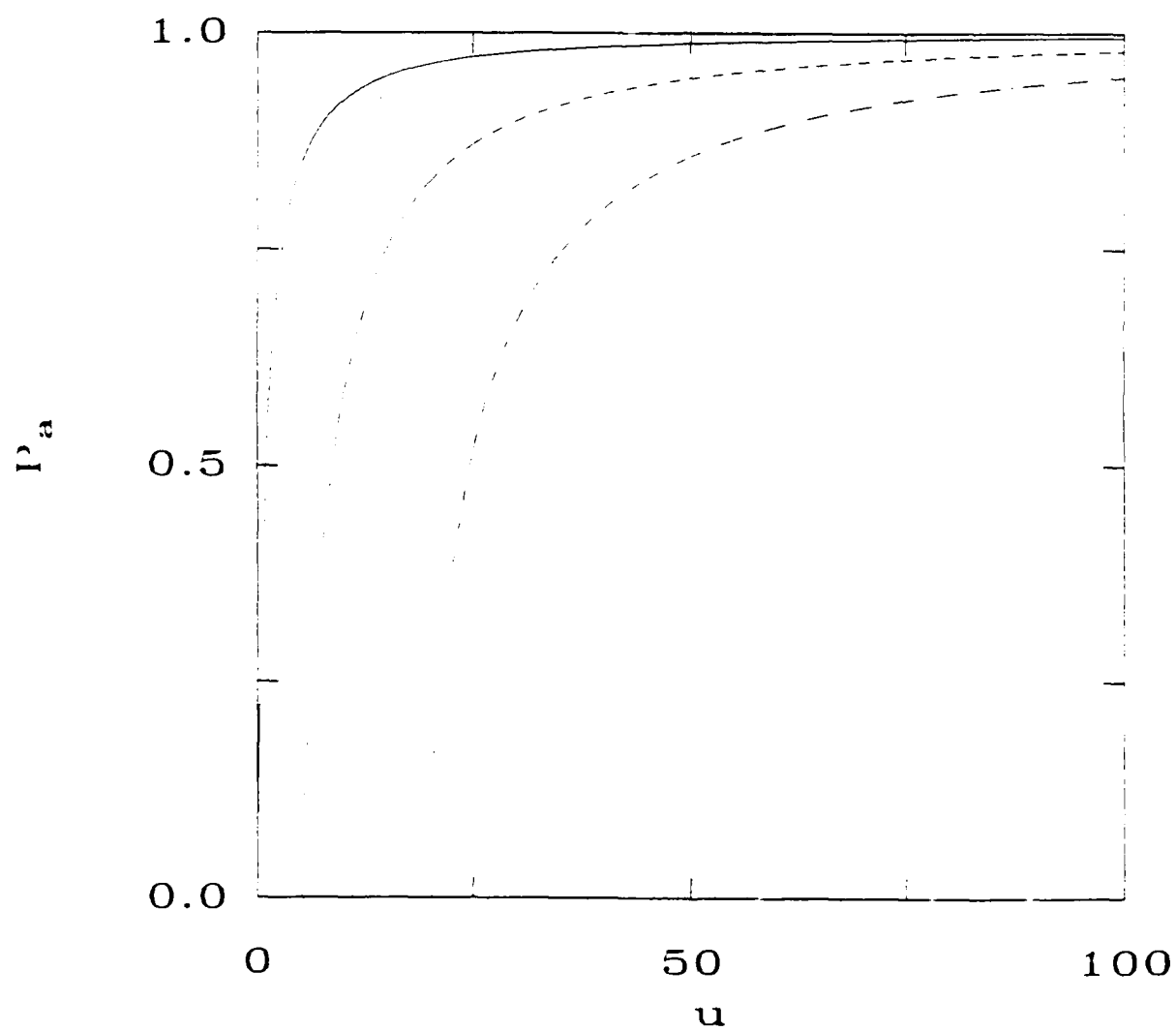


Fig. 2



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